Dynamics of strongly magnetized ejecta in Gamma Ray Bursts

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ABSTRACT

We consider dynamical scales in magnetized GRB outflows, using the solutions to the Riemann problem of expanding arbitrarily magnetized outflows (Lyutikov 2010). For high ejecta magnetization, the behavior of the forward shock closely resembles the so-called thick shell regime of the hydrodynamical expansion. The exception is at small radii, where the motion of the forward shock is determined by the dynamics of *subsonic* relativistic outflows. The behaviors of the reverse shock is different in fluid and magnetized cases: in the latter case, even for medium magnetization, $\sigma \sim 1$, the reverse shock forms at fairly large distances, and may never form in a wind-type external density profile.

1. Introduction

Magnetic fields may play an important dynamical role in the GRB outflows (*e.g.* Lyutikov 2006, 2009). They may power the relativistic outflow through Blandford & Znajek (1977) process (*e.g.* Komissarov 2005), and contribute to particle acceleration in the emission regions. In this paper we discuss the dynamics of the relativistic, strongly magnetized ejecta. The results are based on an exact solution of a one-dimensional Riemann problem of expansion of a cold, strongly magnetized into vacuum and into external medium of density $\rho_{\rm ex}$ (Lyutikov, submitted); they are reviewed in §2.

In application to GRBs, we assumes that the central engine produces jet with density ρ_0 and magnetization σ ($\sigma = B_0^2/\rho_0$; magnetic field is normalized by $\sqrt{4\pi}$), moving with Lorentz factor $\gamma_w \gg 1$. In fact, parameters γ_w and σ are not always independent quantities: at small radii, when the motion of the ejecta is subsonic, they should be determined together with the motion of the boundary, see §3.3. In a supersonic regime, relation between γ_w and σ depends on the details of the flow acceleration (*e.g.* in conical flows we expect $\gamma_w \sim \sqrt{\sigma}$). For generality, we do not assume any relationship between σ and γ_w. The ejecta is moving into external density $\rho_{\rm ex}$.

2. Riemann problem for relativistic expansion of magnetized gas

2.1. Simple waves and forward shock dynamics

Let us assume that the jet plasma is moving with velocity β_w towards the external medium. We found (Lyutikov, submitted) exact self-similar solution of relativistic Riemann problem for the expansion of cold plasma with density ρ_0 and magnetic field B_0 (magnetization parameter $\sigma = B_0^2/\rho_0$; magnetic field is normalized by $\sqrt{4\pi}$), moving initially with velocity v_w towards the vacuum interface

$$\delta_\beta = \delta_\eta^{2/3}\delta_{A,0}^{2/3}\delta_w^{1/3}$$
$$\delta_A = \frac{\delta_{A,0}^{2/3}\delta_w^{1/3}}{\delta_\eta^{1/3}} \tag{1}$$

where the Doppler factors $\delta_a = \sqrt{(1+\beta_a)/(1-\beta_a)}$ are defined in terms of the plasma velocity β, local Alfvén velocity β_A, self-similar parameter $\eta = z/t$, initial wind velocity β_w and the Alfvén velocity in the undisturbed plasma $\beta_{A,0} = \sqrt{\sigma/(1+\sigma)}$. These equations give the velocity β, density $\rho = U_A^2\rho_0/\sigma$ ($U_A = \beta_A/\sqrt{1-\beta_A^2}$) and proper magnetic field, $B = (\rho/\rho_0)B_0$ as a function of the self-similar variable $\eta = z/t$ (expansion of plasma starts at $t = 0, z = 0$ and proceeds into positive direction $z > 0$). We stress that these solutions are exact, no assumptions about the value of the parameter σ and velocity v_w were made.

Particularly simple relations are obtained for plasma initially at rest expanding into vacuum $\beta_w = 0$, $\delta_\beta = 1$ (Lyutikov, submitted). The flow accelerates from rest towards the vacuum interface. The bulk of the flow is moving with Lorentz factor $\gamma' \sim \sigma^{1/3}$. The flow becomes supersonic at $\eta = 0$, at which point $\gamma' = (\sigma/2)^{1/3}$. The vacuum interface moves with Lorentz factor $\gamma'_{vac} = 1 + 2\sigma$. In the observer frame the vacuum interface is moving with $\delta_\eta = \delta_{A,0}^2\delta_w$, which in the limit σ, $\gamma_w \gg 1$ this gives

$$\gamma_{vac} = 4\gamma_w\sigma \tag{2}$$

As the flow expands, the local magnetization

$$\sigma_{loc} = \frac{B^2}{\rho} = \left(\frac{\delta_{A,0}^{2/3}}{\delta_\eta^{1/3}} - \frac{\delta_\eta^{1/3}}{\delta_{A,0}^{2/3}}\right) \tag{3}$$

decreases. At the sonic point $\sigma_{loc} = (\sigma/2)^{2/3}$.

If there is an outside medium with density $\rho_{\rm ex}$, we may identify two expansion regimes. For relativistically strong forward shocks, so that the post-shock pressure is much larger than density, the Lorentz factor of the CD is

$$\gamma_{CD} = \left(\frac{3B_0^2\gamma_w^2}{8\rho_{\rm ex}}\right)^{1/4} \approx \left(\frac{L}{\rho_{\rm ex}c^3}\right)^{1/4} r^{-1/2} \tag{4}$$

(the last approximation assumes $\sigma \gg 1$). For weak forward shocks the velocity of the CD approaches the expansion velocity into vacuum γ_{vac}, Eq. (2). The transition between the relativistically strong and weak shocks occurs for

$$\sigma_{crit} = \left(\frac{3}{2048\gamma_w^2} \frac{\rho_0}{\rho_{\rm ex}} \right)^{1/3} \tag{5}$$

For $\sigma < \sigma_{crit}$, the forward shock is weak.

2.2. Existence of the reverse shock

For cold unmagnetized jets the reverse shock always exists; it is weak for $\gamma_w \leq \sqrt{\rho_0/\rho_{\rm ex}}$ and strong otherwise (Sari & Piran 1995). For magnetized jets the conditions for existence of a reverse shock are more complicated (see also Giannios et al. 2008; Mizuno et al. 2009). There are, in fact, two somewhat different regimes for the existence of a RS in highly magnetized outflows. First, if ejecta is supersonic with respect to the CD (in term of Riemann waves, this transition corresponds to the case when the location of the FS coincides with the location of the rarefaction wave), a strong RS must forms. Secondly, if the ejecta is subsonic with respect to the CD, but moves with velocity higher than the CD, slowing of the ejecta is achieved by a compression wave, which may or may not turn into a reverse shock. One dimensional compression waves are always unstable to shock formation (Landau & Lifshitz 1959). In contrast, multidimensional subsonic outflow need not form shocks. So, formally, the condition for reverse shock is $\gamma_w > \gamma_{CD}$, but in the range $\gamma_{CD} < \gamma_w < 2\gamma_w\sqrt{\sigma}$ the RS may not form, if a more complicated flow patters are allowed. In any case, the RS shock, even if it exists, is weak in this regime.

Conditions for *strong* reverse shock (which implies a highly supersonic flow, with velocity much larger than the Alfvén velocity in the upstream plasma) were derived by Kennel & Coroniti (1984). In the frame of the CD, the reverse shock is moving with (Kennel & Coroniti 1984)

$$\beta_{RS}^{\prime,2} \approx 1 - \frac{1}{\sigma}, \ \ {\rm for} \sigma \gg 1 \tag{6}$$

If $\gamma_w \geq \gamma_{CD}$, the reverse shock is weak (if it exists) and one should use a more detailed calculations of the dynamics of perpendicular shocks of arbitrary strength.

Thus, for $\sigma \geq 1$, the existence of strong RS requires $\gamma_w > 2\gamma_{CD}\sqrt{\sigma}$, which using Eq. (4) gives

$$\gamma_w > \sqrt{6}\sqrt{\frac{\rho_0}{\rho_{\rm ex}}}\sigma^{3/2}, \tag{7}$$

while a weak RS may exist for $\gamma_{CD} < \gamma_w < 2\gamma_w\sqrt{\sigma}$:

$$\sqrt{\frac{3}{8}}\sqrt{\frac{\rho_0}{\rho_{\rm ex}}}\sqrt{\sigma} < \gamma_w < \sqrt{6}\sqrt{\frac{\rho_0}{\rho_{\rm ex}}}\sigma^{3/2}, \tag{8}$$

The Lorentz factor of a strong RS with respect to the contact discontinuity is $\sqrt{\sigma}$ (assuming $\sigma \gg 1$). The Lorentz factor of the reverse shock in the frame of stationary external medium is then

$$\gamma_{RS} = \frac{1}{2}\left(\frac{\sqrt{\sigma}}{\gamma_{CD}} + \frac{\gamma_{CD}}{\sqrt{\sigma}}\right) = \begin{cases} \left(\frac{3}{128}\right)^{1/4} \frac{\sqrt{\gamma_w}}{\sigma^{1/4}} \left(\frac{\rho_0}{\rho_{\rm ex}}\right)^{1/4} & \text{if } \gamma_{CD} \gg \sqrt{\sigma} \\ \frac{\sigma^{1/4}}{6^{1/4}\sqrt{\gamma_w}} \left(\frac{\rho_{\rm ex}}{\rho_0}\right)^{1/4} & \text{if } \gamma_{CD} \ll \sqrt{\sigma} \end{cases} \tag{9}$$

The two cases in Eq. (9) correspond to RS moving in the same direction as the CD, $\gamma_{CD} > \sqrt{\sigma}$, and the RS moving in the opposite direction than the CD, $\gamma_{CD} < \sqrt{\sigma}$. Condition $\gamma_{CD} = \sqrt{\sigma}$ gives

$$\gamma_w = \sqrt{\frac{8}{3}}\sqrt{\frac{\rho_{\rm ex}}{\rho_0}}\sqrt{\sigma} \tag{10}$$

In this case the reverse shock is stationary in the frame of the external medium.

Let us summarize the main results. If the the ratio of ejecta density to external density is $f = \rho_0/\rho_{\rm ex}$, then weak RS can form for $\gamma_w \geq \gamma_{CD}$, which gives $\gamma_w \sim \sqrt{f\sigma}$; strong RS forms for $\gamma_w > \sqrt{\sigma}\gamma_{CD}$, $\gamma_w \geq \sqrt{f\sigma^3}$. RS shock propagates in the forward direction for $\gamma_{CD} > \sqrt{\sigma}$, $\gamma_w > \sqrt{\sigma/f}$. Forwards shock is relativistically weak for $\gamma_{CD} \geq \sigma$, $\gamma_w > \sigma^{3/2}/\sqrt{f}$ and becomes non-relativistic for $\gamma_{CD} \sim 1$, $\gamma_w < 1/\sqrt{\sigma f}$.

3. Dynamics of magnetized flows in GRBs

In this section we apply the previous relationships to consider dynamics of magnetized flows in GRBs, generalizing discussion of Sari & Piran (1995) to strongly magnetized flows. We will derive main results in a thin shell approximation (not to be confused with a thin shell case, see below), assuming that the distances between the forward shock, the contact discontinuity and the reverse shock are small. The velocity of the shocks and contact discontinuity are determined from the local force balance conditions. More precisely, they are determined by the *local* solutions to the Riemann problem of the decay of the discontinuity of the flow: there is no memory in the flow. Thin shell approximation is likely to be applicable, since for reasonable GRB parameters the reverse shock never stalls while expansion is relativistic, see discussion after Eq. (20).

The ejecta flow is taken to expands conically and carrying toroidal magnetic field. We assume that the central source operates for time $\Delta t_s = \Delta/c$ (Δ is the initial width of the launched shell) and produces a wind with magnetization $\sigma \gg 1$ (magnetization $\sigma = B^2/\rho$ is twice the ratio of magnetic to particle energy in plasma frame; magnetic field is normalized by $\sqrt{4\pi}$), moving with the Lorentz factor γ_w. For spherical expansion (expansion along conical surfaces) of magnetized flows into vacuum, the magnetization parameter σ remains constant outside the fast magnetosonic surface (Michel 1973; Vlahakis & Königl 2003). As we will see, the above assumption (that the central source produces a flow with a given γ_w and σ) is not self-consistent for small radii, where the reverse shock does not form. In this case of subsonic expansion, the flow dynamics cannot be specified *ad hoc*: it needs to be determined self-consistently with the motion of the boundaries.

The wind luminosity is assumed to be $L_{\rm iso} = E_{\rm iso}/\Delta t_s$ where $E_{\rm iso}$ is the isotropic equivalent energy released by the central source. Luminosity is produced in a form of Poynting and particle fluxes

$$L = 4\pi r^2 \gamma_w^2 (B_0^2 + \rho_0) = 4\pi r^2 \gamma_w^2 B_0^2 \frac{1+\sigma}{\sigma} \tag{11}$$

We are interested in the case $\sigma \geq 1$. For numerical estimates we will use the typical values for long GRBs: $L_{\rm iso} = 10^{51}$ erg s^{-1}, $\Delta t_s = 100$ s, $E_{\rm iso} = 10^{53}$ erg, $\gamma_w = 300$. External density is $\rho_{\rm ex} = m_p n$.

3.1. Forward shock dynamics

In case of magnetized ejecta, as well as in the hydrodynamical case (Sari & Piran 1995), the important scales in the problem (Sedov scale l_S (13), energy scale r_E (14), reverse shock formation scale r_N (18), reverse shock crossing scale r_Δ (21) and spreading distance r_S (22)) are related by a quantity (Sari & Piran 1995)

$$\xi = \sqrt{\frac{l_S}{\Delta}} \gamma_w^{-4/3}, \ \ r_N/\xi = r_E = \sqrt{\xi} r_\Delta = \xi^2 r_s \tag{12}$$

In the hydrodynamical case, the parameter ξ determines whether the reverse shock and the rarefaction wave reach the whole ejecta before most of the energy is transferred to the forward shock, $\xi > 1$, or later, $\xi < 1$. The dynamics of magnetized ejecta generally follows the hydrodynamic thick case, though the meaning of some radii change (*e.g.*, in case of strongly magnetized ejecta r_N is the scale of RS formation).

There is a number of typical radii where dynamics of the outflow changes. There is Sedov radius

$$l_S \sim \left(\frac{E_{\rm iso}}{\rho_{\rm ex} c^2} \right)^{1/3} = 4 \times 10^{18}\, {\rm cm}\, n^{-1/3} \tag{13}$$

where the ejecta and the swept-up ISM material become non-relativistic.

There is radius r_E, where the ejecta deposits approximately half of the energy or momentum to the external medium. For supersonic flows, which reached terminal Lorentz factor γ_w, equating energy in the shocked medium $\gamma_w^2 \rho_{\rm ex} c^2 r_E^3$ to the total energy $E_{\rm iso}$, gives (Rees & Meszaros 1992)

$$r_E \sim \left(\frac{E_{\rm iso}}{\gamma_w^2 \rho_{\rm ex} c^2} \right)^{1/3} = \frac{l_S}{\gamma_w^{2/3}} = 9 \times 10^{16}\, {\rm cm}\, n^{-1/3} \tag{14}$$

r_E depends exclusively on the total energy of the explosion and not on its form (magnetic or baryonic). For radii smaller than r_E the ejecta's and the forward shocks' Lorentz factors remain constant and equal to the initial Lorentz factor γ_w. For larger radii the flow enters the self-similar Sedov-Blandford-McKee stage, with Lorentz factor decreasing according to

$$\gamma = \left(\frac{l_s}{r} \right)^{3/2} \tag{15}$$

In case of pure baryonic flow, and only in that case, r_E is also the radius when the swept-up mass equals the ejecta mass divided by γ_w

$$r_M \sim \left(\frac{M_0}{\gamma_w \rho_{\rm ex} c^2}\right)^{1/3} = \left(\frac{E_K}{\gamma_w^2 \rho_{\rm ex} c^2}\right)^{1/3} \tag{16}$$

Here $E_K = E_{\rm iso}/(1+\sigma)$ is the energy associated with bulk motion of matter. Only in the case of zero magnetization r_E equals r_M, since in that case $E_{\rm iso} = E_K = M_0 \gamma_w$. [1] For highly magnetized outflow $\frac{r_M}{r_E} = \sigma^{-1/3} \ll 1$.

The above description of the forward shock dynamics is, in fact, applicable only in the so called thin shell case, $\xi > 1$ (see Eq. (12). In this case the reverse shock quickly crosses the ejecta, which becomes causally connected so that all of the ejecta interacts with the external medium. Alternatively, in the thick shell case, $\xi < 1$ (see Eq. (12), the reverse shock does not have time to cross the ejecta before the causally connected shocked part starts to decelerate at smaller radius r_N, Eq. (18). The Lorentz factor starts decreasing, but since new material and new momentum is being added to the shocked part of the ejecta, the ejecta and the forward shock behave effectively as a self-similar shock with energy supply

$$\gamma = \left(\frac{L}{\rho_{\rm ex} c^2}\right)^{1/4} \frac{1}{\sqrt{r}} = \frac{l_s^{3/4}}{\Delta^{1/4}\sqrt{r}} \tag{17}$$

Since in the thick shell case the Lorentz factor starts to decelerate earlier than in the thin shell case, the rate of energy transfer to the external medium is smaller, so that the ejecta gives most of its initial energy to the ISM at larger distances $r_\Delta > r_E$ (Sari & Piran 1995).

3.2. Formation and dynamics of the reverse shock

The weak reverse shock may form at (see Eq. (8))

$$r_N = \frac{1}{\gamma_w^2}\sqrt{\frac{3L}{2\pi \rho_{ISM} c^3}} \approx \frac{1}{\gamma_w^2}\frac{l_S^{3/2}}{\sqrt{\Delta}} = 10^{16}\,{\rm cm}\, n^{-1/2} \tag{18}$$

RS becomes strong at $r_{RS,\rm strong} \sim \sigma r_N$ (see Eq. (7)).

If the outside medium is stellar wind, strong RS forms immediately if

$$\gamma_w > \left(\frac{3}{2\pi}\frac{L v_{wind}\sigma^2}{c^3 \dot{M}}\right)^{1/4} = 220 \sigma^{1/2} L_{51}^{1/4} v_{wind,8}^{1/4} \left(\frac{\dot{M}}{10^{-8} M_\odot/{\rm yr}}\right)^{-1/4} \tag{19}$$

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[1]The two radii r_M and r_E were confused by Zhang & Kobayashi (2005), who "define the deceleration radius using E_K alone [] where the fireball collects $1/\gamma_w$ of fireball rest mass". According to Zhang & Kobayashi (2005) "only the kinetic energy of the baryonic component (E_K) defines the afterglow level", while magnetic energy is transferred at unspecified "later" time. This is incorrect (Lyutikov 2005).

where $v_{wind,8}$ is the velocity of the progenitors wind in thousands kilometers per second. For smaller γ_w, no RS forms ever (for weak shocks, one should put $\sigma \to 1$).

The reverse shock could stall (in the observer frame) at (assuming $\sigma \gg 1$)

$$r_{RS,stall} = \frac{1}{\sigma}\sqrt{\frac{3L}{8\rho_{\rm ex}c^3}} \sim \frac{1}{\sigma}\frac{l_S^{3/2}}{\sqrt{\Delta}} = 3\times 10^{21}{\rm cm} n^{-1/2}\sigma^{-1} \tag{20}$$

Since $r_{RS,stall}$ is typically larger than l_S, RS does not stall during the relativistic expansion phase; thus, the thin shell approximation is generally applicable.

In the unmagnetized case, an important quantity is the radius when the RS crosses the ejecta

$$r_\Delta \sim \left(\frac{E\Delta}{\rho_{\rm ex}c^2}\right)^{1/4} = l_S^{3/4}\Delta^{1/4} = 10^{17}\,{\rm cm}\,n^{-1/4} \tag{21}$$

In the magnetized case, r_Δ is still a good approximation for the RS crossing radius, but with two cavities. First, a delayed onset of the RS, see (18), delays the RS crossing moment. Since $r_N/r_\Delta = \xi^{3/2}$, this delay is not important for $\xi < 1$ (the thick shell case, generally applicable to the magnetized ejecta). Also, for subsonic outflows (see below), r_Δ is the distance where the back of the outflow catches with the CD, see §3.3.

Second, magnetized shell is necessarily expanding, so that the tail part of the flow is moving with $\gamma \sim \gamma_{CD}/(2\sqrt{\sigma})$. The typical shell spreading distance is

$$r_S \sim \Delta\gamma_w^2 = 3\times 10^{17}{\rm cm} \tag{22}$$

Since the spreading occurs with Alfven velocity, the tail part of the flow catches with the CD in the Blandford-McKee phase at $r_{tail} \sim \sqrt{\sigma} r_\Delta$. We stress that spreading of magnetic shell, unlike of the cold baryonic shell, is unavoidable consequence of the high internal pressure (spreading of cold baryonic shell requires internal motion). This is the reason why magnetic outflows are similar to the thick shell case of the baryonic outflows.

3.3. Dynamics of subsonic expansion

Thus, for a given γ_w and σ, at distances $r < r_N$ a rarefaction wave is launched into the flow, while the flow accelerates to Lorentz factors larger than γ_w. This implies that our assumption that a flow has a given γ_w and σ is not justified at $r < r_N$: at these distances the flow is effectively subsonic and its dynamics needs to be solved self-consistently, taking into account interaction with the external medium. The subsonic outflow may be considered as a collection of outgoing fast magnetosonic waves propagating from the central source, which constantly re-energize the FS. The outflow may be separated into two stages, which we will call "early" and "late", depending on whether or not most of the fast waves emitted by the central source have caught up with the CD and their energy has been given to the circumburst medium. The transition between two stages

occurs at the moment, which is similar to the shell crossing radius in the supersonic case, except that in the case of subsonic expansion the CD is decelerating all the time, but with different laws before and after the transition.

At the "early" stage the CD is constantly re-energized by the fast-magnetosonic waves propagating from the central source. The motion of the CD is determined by the luminosity at the retarded time t':

$$L_\Omega(t') \sim \rho_{\rm ex} c^3 \gamma_{CD}^4 r^2 \tag{23}$$

(this is a condition of pressure balance between the wind and the ram pressure of ISM in the frame of the CD). For constant luminosity Eq. (23) gives Eq. (17) for the Lorentz factor of the CD. This is exactly the same estimate as for the intermediate scale in $\xi < 1$ supersonic flows, since $L_\Omega \sim E_{\rm iso} c/\Delta$; also this is the same scaling as in the case of relativistic fluid reverse shock (Sari 1997). We stress that in the limit of strong FS, Eq. (5), *the FS dynamics is independent of the composition of the flow*, only the total power is important, Eq. (23).

The early stage lasts for $r < r_\Delta$, Eq. (21). At larger radii the flow enters the self-similar Sedov-Blandford-McKee stage. At this stage, only a fraction of the shell is interacting with the external medium, while the newly shocked ejecta material keeps adding energy and momentum to the shocked shell and the ISM, which evolve, effectively, as a flow with energy supply. Finally, for $r > r_\Delta$ the shock enters Blandford-McKee stage, with Lorentz factor given by Eq. (15).

4. Discussion

In this paper we discuss the dynamics of strongly magnetized outflows in GRBs. We find that the evolution of the forward shock driven by strongly magnetized outflows are qualitatively the same as in the case of fluid shocks. The definitions of radii r_N, r_E and r_Δ involve only the total energy of the ejecta, it's thickness and initial Lorentz factor, and *not* the information about it's content, *e.g.*, parameter σ. The typical radii (12) are the same for two flows (*cf.* Eq. (12) of the present paper and Eq. (9-10) of Sari & Piran 1995). These similarities may be understood, first, by noting that jump conditions in perpendicular magnetized shocks may be reduced to fluid shock jump conditions, with an appropriate choice of the equation of state, and, second, by the fact that the thin shell approximation is applicable in our case (so that the global conservation of the toroidal magnetic flux, which modifies the global flow dynamics (Kennel & Coroniti 1984), is not important). Another reason for this similarity is that magnetic field behaves in many respects as a fluid with internal pressure. The only difference in the dynamics of the forward shocks driven by magnetized and fluid flows occurs for supersonic flows, $\gamma_w > \sqrt{\sigma}\gamma_{CD}$, at very early stages $r \leq r_N$ or $r \leq r_E$, see Fig. (1). Qualitatively, magnetized outflows are similar to thick shell hydrodynamic outflow, $\xi < 1$ at $r > r_N$.

Only at very early times, at $r < r_N$, the forward shock bears information about anergy content: forward shock is coasting with γ_w =const in the fluid case and decelerating $\gamma \propto r^{-1/2}$ in

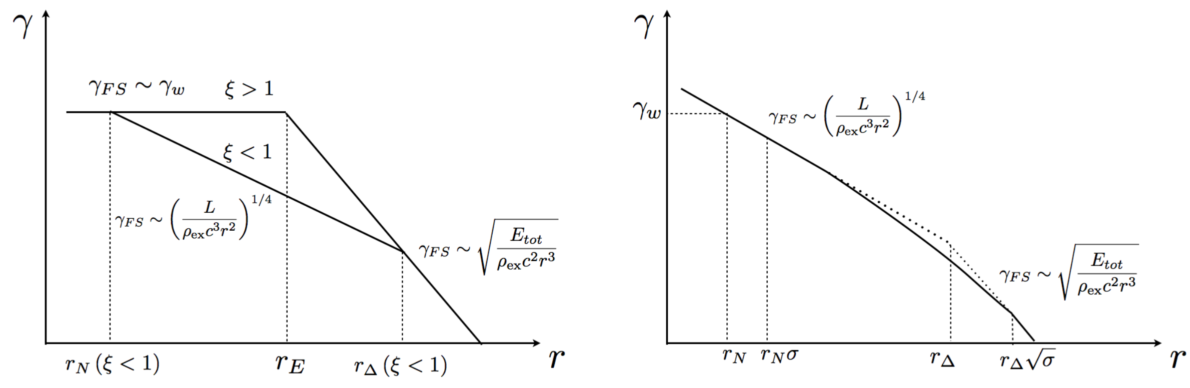

Fig. 1.— Evolution of the Lorentz factor of the forward shock for matter-dominated (Left panel) and Poynting flux-dominated models (Right panel). Matter-dominated ejecta coasts with the injection Lorentz factor γ_w until either r_E (for $\xi > 1$, thin shell case) or until r_N (for $\xi < 1$, thick shell case). At r_N reverse shock becomes strong. At large radii ($r > r_E$ or $r > r_\Delta$) the outflow enters the Sedov-Blandford-McKee regime. For highly magnetized ejecta the Lorentz factor of the CD and the FS initially decreases $\Gamma_{ISM} \propto r^{-1/2}$, changing to Sedov-Blandford-McKee regime $\Gamma_{ISM} \propto r^{-3/2}$ approximately at r_Δ. Reverse shock is launched at r_N and becomes strong at $r_N\sigma$. Due to internal expansion of the magnetized shell, the back of the shell catches with the CD at distance $\sim r_\Delta\sqrt{\sigma}$. This is the reason why at distances close to r_Δ the Lorentz factor starts decreasing below the $r^{-1/2}$ law.

the magnetized case. Dynamics of the reverse shock is quite different in case of high magnetization. First, the reverse shock forms at a finite distance from the source (Eq. 18), and may not form at all in a wind environment, (Eq. 19). This fact may be related to observed paucity of optical flashes in the *Swift* era (Gomboc et al. 2009). (The standard model had a clear prediction, of a bright optical flare with a definite decay properties (Sari et al. 1996; Meszaros & Rees 1997). Though a flare closely resembling the predictions was indeed observed (GRB990123, Mészáros & Rees 1999), this was an exception.)

In addition, at distances $r_N < r < \sigma r_N$, where the RS is weak, the formation of the RS shock depends on the details of the flow: RS forms if the flow is strictly radial, but need not to form if the the flow pattern is more complicated. We suggest that optical variability often seen in GRBs (*e.g.* GRB021004 and most notoriously GRB080916C) is a reflection of the non-trivial flow patterns and the corresponding non-steady RS formation. Also, a recent detection of high polarization in optical (Steele et al. 2009) indicates a presence of an ordered magnetic field in the ejecta.

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